

RECENT THEORETICAL DEVELOPMENTS IN $\bar{B} \rightarrow X_s \ell^+ \ell^-$ DECAYS^a

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We present a concise review of the theoretical status of the rare semileptonic $\bar{B} \rightarrow X_s \ell^+ \ell^-$ decays in the standard model. Particular attention is thereby devoted to the recent theoretical progress concerning, on the one hand the next-to-next-to-leading order QCD calculation and, on the other hand the analysis of phenomenological important subleading electroweak effects.

1 Introduction

The rare semileptonic $b \rightarrow s \ell^+ \ell^-$ transitions have been observed for the first time by Belle and BaBar in form of the exclusive $\bar{B} \rightarrow K^{(*)} \ell^+ \ell^-$ modes.^{1,2,3} Recently also inclusive measurements have become available.^{2,3,4} Like all other flavor-changing-neutral-current processes, these channels are important probes of short-distance physics. Their study can yield useful complementary information, when confronted with the less rare $b \rightarrow s \gamma$ decays, in testing the flavor sector of the Standard Model (SM). In particular, a precise measurement of the inclusive $\bar{B} \rightarrow X_s \ell^+ \ell^-$ mode would be welcome in view of New Physics (NP) searches, because it is amenable to a clean theoretical description for dilepton invariant masses, $m_{\ell\ell}^2 \equiv q^2$, below and above the $c\bar{c}$ resonances, namely in the ranges $1 \text{ GeV}^2 \leq q^2 \leq 6 \text{ GeV}^2$ and $q^2 \geq 14.4 \text{ GeV}^2$.⁵ Each window has its own assets and drawbacks, related on the one hand to experimental issues, such as event rates, identification and detection efficiencies, on the other hand to theoretical questions, concerning the significance of parametric errors, perturbative and non-perturbative effects. In the following we present results that will cover both regions, performing a thorough study of the associated uncertainties. Unfortunately, however, it is not possible to compare these predictions with experiment, as it is still too early to fit a q^2 distribution to the data. To allow the comparison with

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the existing inclusive measurements, we thus reexamine the SM prediction for the $\bar{B} \rightarrow X_s \ell^+ \ell^-$ branching ratio, $\text{BR}_{\ell\ell}$, including all known QCD and electroweak effects. Finally, we update the SM result for the position of the zero of the Forward-Backward (FB) asymmetry, q_0^2 .

2 Theoretical Framework

The calculation of the partonic decay rate of $\bar{B} \rightarrow X_s \ell^+ \ell^-$ consists of several parts that are worth recalling. Perturbative QCD effects play an important role, due to the presence of large logarithms of the form $L \equiv \ln m_b/M_W$, that can be resummed using the machinery of Operator Product Expansion (OPE) and Renormalization Group (RG) improved perturbation theory. Factoring out the Fermi constant G_μ and the electromagnetic coupling α , the amplitude receives contributions of $O(\alpha_s^n L^{n+1})$ at the Leading Order (LO), of $O(\alpha_s^n L^n)$ at the Next-to-Leading Order (NLO), and of $O(\alpha_s^n L^{n-1})$ at the Next-to-Next-to-Leading Order (NNLO) in QCD.

To achieve the necessary resummation, one works in the framework of an effective low-energy theory with five active quarks, three active leptons, photons and gluons, obtained by integrating out heavy degrees of freedom characterized by a mass scale $M \geq M_W$. At LO in the OPE the effective on-shell Lagrangian relevant for the $b \rightarrow s \ell^+ \ell^-$ transition at a scale μ is given by

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{QCD} \times \text{QED}} + \frac{4G_\mu}{\sqrt{2}} V_{ts}^* V_{tb} \sum_i C_i(\mu) Q_i. \quad (1)$$

Here the first term is the conventional QCD–QED Lagrangian for the light SM particles. In the second term V_{ij} denotes the elements of the Cabibbo-Kobayashi-Maskawa matrix and $C_i(\mu)$ are the Wilson coefficients of the corresponding operators Q_i built out of the light fields.

Neglecting for the moment subleading electroweak effects, as well as the QCD penguin operators Q_3 – Q_6 , that are suppressed by small Wilson coefficients, the remaining physical operators arising in the SM can be written as

$$\begin{aligned} Q_1 &= (\bar{s}_L \gamma_\mu T^a c_L)(\bar{c}_L \gamma^\mu T^a b_L), & Q_2 &= (\bar{s}_L \gamma_\mu c_L)(\bar{c}_L \gamma^\mu b_L), \\ Q_7^\gamma &= \frac{e}{g^2} m_b (\bar{s}_L \sigma^{\mu\nu} b_R) F_{\mu\nu}, & Q_8^g &= \frac{1}{g} m_b (\bar{s}_L \sigma^{\mu\nu} T^a b_R) G_{\mu\nu}^a, \\ Q_9 &= \frac{e^2}{g^2} (\bar{s}_L \gamma_\mu b_L) \sum_\ell (\bar{\ell} \gamma^\mu \ell), & Q_{10} &= \frac{e^2}{g^2} (\bar{s}_L \gamma_\mu b_L) \sum_\ell (\bar{\ell} \gamma^\mu \gamma_5 \ell), \end{aligned} \quad (2)$$

where the sum over ℓ extends over all lepton fields, e (g) is the electromagnetic (strong) coupling constant, q_L and q_R are the chiral quark fields, $F_{\mu\nu}$ ($G_{\mu\nu}^a$) is the electromagnetic (gluonic) field strength tensor, and T^a are the color matrices, normalized so that $\text{Tr}(T^a T^b) = \delta^{ab}/2$.

Unlike the case of $b \rightarrow s\gamma$, the $b \rightarrow s \ell^+ \ell^-$ amplitude involves large logarithms even in the absence of QCD interactions, since the current-current operators Q_1 and Q_2 , as well as Q_3 – Q_6 mix into the vector-like semileptonic operator Q_9 at the one-loop level. QCD contributions from the magnetic operators Q_7^γ and Q_8^g , and the axial-vector-like semileptonic operator Q_{10} enter formally at the NLO, but turn out to be numerically an $O(1)$ correction to the LO result. In order to gain an accuracy below the 10% level on the $b \rightarrow s \ell^+ \ell^-$ decay rate a complete resummation of NNLO QCD logarithms has thus to be performed. Contrary thereto, a similar precision is achieved for $b \rightarrow s\gamma$ already at the NLO,^{6,7} accentuating once again the difference between rare and radiative modes in view of the RG improved perturbation theory.

3 Recent Perturbative Standard Model Calculations

The aforementioned NNLO QCD computation has required the computation of *i*) the $O(\alpha_s)$ corrections to the relevant Wilson coefficients,⁸ *ii*) the $O(\alpha_s)$ contributions to the associated

matrix elements,^{5,9,10,11,12,13} and *iii*) the $O(\alpha_s^2)$ Anomalous Dimension Matrix (ADM) describing the mixing of physical dimension-five and six operators.^{14,15} Nearly all the ingredients of the NNLO QCD calculation involve a considerable degree of technical sophistication and have been performed independently by at least two groups, sometimes using different methods. However, the most complex part of the whole enterprise, the calculation of the three-loop $O(\alpha_s^2)$ ADM describing the mixing of Q_1 – Q_6 into Q_1 – Q_9 has been completed only very recently.¹⁵ Although the $O(\alpha_s^2)$ matrix elements of Q_3 – Q_6 have not been calculated so far,^b the NNLO QCD computation of $b \rightarrow s\ell^+\ell^-$ can be called practically complete, as the $O(\alpha_s^2)$ matrix elements of Q_1 and Q_2 ,⁵ as well as Q_9 ,¹³ evaluated for arbitrary q^2 , are now also available. In their sum the numerical impact of the NNLO QCD corrections on the differential decay rate amounts to around -20% (-25%) in the low- q^2 (high- q^2) region, and leads to a reduction of the renormalization scale uncertainties from around $\pm 20\%$ ($\pm 15\%$) to $\pm 5\%$ ($\pm 3\%$). In the case of the FB asymmetry, NNLO QCD corrections are not less important, as they shift q_0^2 by around $+15\%$ and reduce the renormalization scale dependences from around $\pm 20\%$ to $\pm 3\%$.

In contrast to the track record of QCD corrections, the possible importance of subleading electroweak effects in $b \rightarrow s\ell^+\ell^-$ has been realized only quite recently. As shown in the case of radiative decays,¹⁷ they may be as important as the higher order QCD effects. $\text{BR}_{\ell\ell}$ is generally parameterized in terms of α , but the scale at which α should be evaluated is, in principle, undetermined until higher order electroweak effects are taken into account. This has led most authors to use $\alpha(m_{\ell\ell}) \approx \alpha(m_b) \approx 1/133$. Indeed, in the absence of an $O(\alpha)$ calculation, there is no reason to consider $\alpha(m_{\ell\ell})$ more appropriate than, say, $\alpha(M_W) \approx 1/128$. As $\text{BR}_{\ell\ell}$ is proportional to α^2 , the ensuing uncertainty of almost 8% is, compared to the precision achieved in the QCD calculation, not at all negligible.^{13,18} This uncertainty, which has affected previous analyses, has been recently reduced to 2% , by a calculation of the numerical dominant subleading electroweak effects,¹³ showing that after the inclusion of the latter corrections the appropriate prefactor in $\text{BR}_{\ell\ell}$ is in fact $\alpha(m_b)$. Due to accidental cancellation, the impact of the calculated electroweak effects is however rather limited: it amounts to almost -2% in $\text{BR}_{\ell\ell}$, whereas q_0^2 is changed by around $+2\%$, and leaves the scale uncertainties practically unchanged.

4 Phenomenology

In order to cancel the strong m_b -dependence of the differential $\bar{B} \rightarrow X_s\ell^+\ell^-$ decay rate, it is customary to normalize the latter to the experimental value of the Branching Ratio (BR) for the inclusive semileptonic decay $\text{BR}[\bar{B} \rightarrow X_c\ell\nu]$. However, this normalization introduces a strong m_c -dependence, which is not known very accurately. The ensuing uncertainty is about 8% .⁵ An alternative procedure^{6,18} consists in normalizing the $\bar{B} \rightarrow X_s\ell^+\ell^-$ decay width to $\Gamma[\bar{B} \rightarrow X_u\ell\nu]$, and then to express $\text{BR}[\bar{B} \rightarrow X_u\ell\nu]$ in terms of $\text{BR}[\bar{B} \rightarrow X_c\ell\nu]$ and of the ratio

$$C = \left| \frac{V_{ub}}{V_{cb}} \right|^2 \frac{\Gamma[\bar{B} \rightarrow X_c\ell\nu]}{\Gamma[\bar{B} \rightarrow X_u\ell\nu]}, \quad (3)$$

which can be computed with better accuracy.⁶ In other words, defining $\hat{s} \equiv q^2/m_b^2$, one can write the normalized differential decay rate as

$$R(\hat{s}) = \frac{\text{BR}[\bar{B} \rightarrow X_c\ell\nu]}{C} \left| \frac{V_{ub}}{V_{cb}} \right|^2 \frac{1}{\Gamma[\bar{B} \rightarrow X_u\ell\nu]} \frac{d\Gamma[\bar{B} \rightarrow X_s\ell^+\ell^-]}{d\hat{s}}, \quad (4)$$

where

$$\frac{d\Gamma[\bar{B} \rightarrow X_s\ell^+\ell^-]}{d\hat{s}} = \frac{G_\mu^2 m_{b,\text{pole}}^5 |V_{ts}^* V_{tb}|^2}{48\pi^3} \left(\frac{\alpha(m_b)}{4\pi} \right)^2 (1 - \hat{s})^2 \left\{ \left(4 + \frac{8}{\hat{s}} \right) \left| \tilde{C}_{7,R}^{\text{eff}}(\hat{s}) \right|^2 \right.$$

^bIn $\bar{B} \rightarrow X_s\gamma$ the $O(\alpha_s^2)$ matrix elements of Q_3 – Q_6 reduce the branching ratio by around 1% .¹⁶

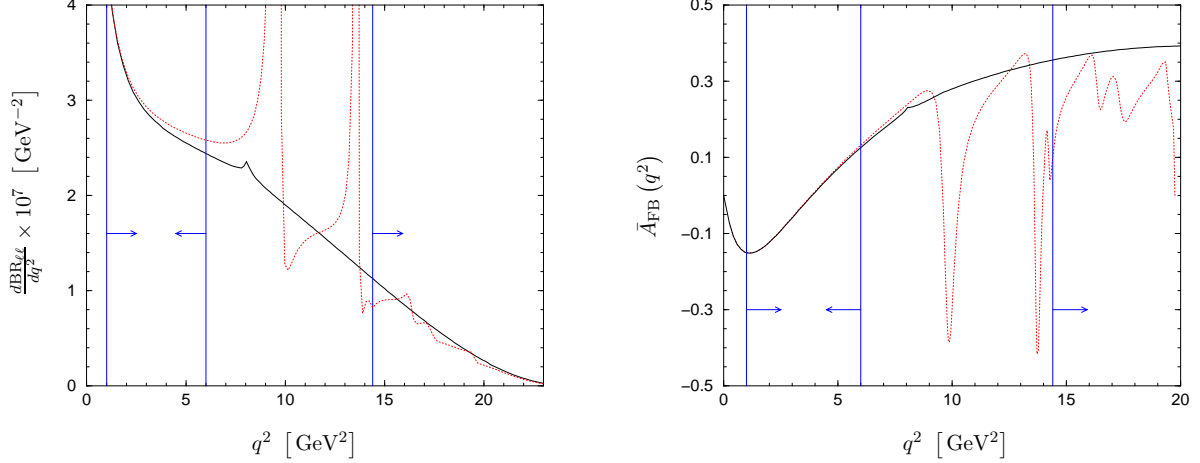


Figure 1: NNLO QCD predictions of $d\text{BR}_{\ell\ell}/dq^2$ (left) and of $\bar{A}_{\text{FB}}(q^2)$ (right) with (dotted red line) and without (solid black line) a guesstimate of $c\bar{c}$ effects. The solid blue lines indicate the boundaries of the reference regions.

$$+(1+2\hat{s})\left(\left|\tilde{C}_{9,R}^{\text{eff}}(\hat{s})\right|^2+\left|\tilde{C}_{10,R}^{\text{eff}}(\hat{s})\right|^2\right)+12\text{Re}\left(\tilde{C}_{7,R}^{\text{eff}}(\hat{s})\tilde{C}_{9,R}^{\text{eff}}(\hat{s})^*\right)+\frac{d\Gamma^{\text{Brems}}}{d\hat{s}}\Bigg\}, \quad (5)$$

and the effective Wilson coefficients carrying the label R , include all real and virtual $O(\alpha_s)$ corrections,^{5,9,10} whereas the last term denotes the finite $O(\alpha_s)$ bremsstrahlungs corrections.¹¹

The use of the b -quark pole mass in Eq. (5) and in the calculation of the semileptonic width leads to large perturbative QCD corrections both in the numerator and denominator of Eq. (4). Since this is an artifact of the choice of scheme for the b -quark mass, they clearly tend to cancel in the ratio. As a second improvement with respect to previous analyses, we keep terms through $O(\alpha_s^2)$ in the denominator and expand the ratio in Eq. (4) in powers of α_s . By making explicit the cancellation of large contributions, the convergence and stability of the perturbative series improves, as we have verified explicitly. Like in previous analyses, due to the peculiarity of the perturbative expansion for $b \rightarrow s\ell^+\ell^-$, we retain some large scheme-independent higher order terms in the amplitude squared: in particular, all terms in Eq. (5) that are quadratic in the effective Wilson coefficients are expanded in α_s and terms up to $O(\alpha_s)$ are kept. A careful study shows,¹³ that with the two aforementioned improvements the residual scale dependence is about 50% smaller than in all preceding analyses.

After including power corrections of $O(1/m_b^2)$ ¹⁹ and $O(1/m_c^2)$ ²⁰ to the differential decay rate of $\bar{B} \rightarrow X_s\ell^+\ell^-$, as well as $O(1/m_b^2)$ corrections to $\Gamma[\bar{B} \rightarrow X_c\ell\nu]$ and $\Gamma[\bar{B} \rightarrow X_u\ell\nu]$, and expanding Eq. (5) in inverse powers of m_b and m_c , we obtain for the normalized differential rate integrated over the low- q^2 region

$$\text{BR}_{\ell\ell}(1\text{ GeV}^2 \leq q^2 \leq 6\text{ GeV}^2) = \left(1.57_{-0.10}^{+0.11}|_{m_t} +0.07_{-0.07}^{+0.07}|_{m_b} +0.06_{-0.07}^{+0.06}|_{\text{scale}} +0.05_{-0.05}^{+0.05}|_{c\bar{c}} +0.05|_C\right) \times 10^{-6}. \quad (6)$$

Including other subleading parametric uncertainties, the total error is about 10% and is dominated by the uncertainty on the t -quark mass, which will soon be reduced by a factor of two by CDF and D0. Moreover, the substantial error from the b -quark mass is an artifact of the employed scheme, which could be reduced drastically by changing the latter. Hence it should not be interpreted as a limitation.

Similarly, we find for the high- q^2 window

$$\text{BR}_{\ell\ell}(q^2 \geq 14.4\text{ GeV}^2) = \left(4.02_{-0.71}^{+0.71}|_{m_b} +0.24_{-0.23}^{+0.24}|_{m_t} +0.13_{-0.13}^{+0.13}|_{\text{scale}} +0.13_{-0.13}^{+0.13}|_{c\bar{c}} +0.12|_C\right) \times 10^{-7}, \quad (7)$$

which is dominated by an error of around 15% related to the kinematical cut on q^2 .⁵ In fact, even though Eq. (5) is not explicitly affected by $O(1/m_b)$ corrections, the physical observable

is sensitive to $O(1/m_b)$ terms, that arise from the well-known relation between the mass of the B -meson and the b -quark in heavy quark effective theory.

Finally, our prediction for the integral of the partonic decay rate over the full spectrum reads

$$\text{BR}_{\ell\ell}(q^2 \geq 4m_\mu^2) = (4.58 \pm 0.18_{\text{scale}} \pm 0.66_{\text{para}}) \times 10^{-6}, \quad (8)$$

where the parametric error takes into account, besides the power corrections of $O(1/m_b^2)$ and $O(1/m_c^2)$, a guesstimate of the uncertainty related to higher $c\bar{c}$ resonances, evaluated by means of experimental data on $e^+e^- \rightarrow X_c$ using a dispersion relation.²¹ The numerical size of the latter correction can be easily assessed from the left plot of Figure 1,^c showing the NNLO QCD prediction for the differential branching ratio as a function of q^2 .

Unfortunately for all of us who hoped to discover NP in rare $b \rightarrow s$ transitions, the estimate in Eq. (8) compares fairly good with the recent experimental world average²²

$$\text{BR}_{\ell\ell}^{\text{exp}} = \left(6.2_{-1.1}^{+1.1} |_{\text{stat}} {}_{-1.3}^{+1.6} |_{\text{syst}}\right) \times 10^{-6}. \quad (9)$$

Even worse, it agrees amazingly well with the preliminary result by Belle,² which suggests a somewhat lower central value for the combination of all inclusive measurements.

Facing such bad news, let's turn our attention to the FB asymmetry, which in its so-called normalized form is defined as

$$\bar{A}_{\text{FB}}(\hat{s}) = \frac{1}{d\Gamma[\bar{B} \rightarrow X_s \ell^+ \ell^-]/d\hat{s}} \int_{-1}^1 d\cos\theta_\ell \frac{d^2\Gamma[\bar{B} \rightarrow X_s \ell^+ \ell^-]}{d\hat{s} d\cos\theta_\ell} \text{sgn}(\cos\theta_\ell), \quad (10)$$

where

$$\begin{aligned} \int_{-1}^1 d\cos\theta_\ell \frac{d^2\Gamma[\bar{B} \rightarrow X_s \ell^+ \ell^-]}{d\hat{s} d\cos\theta_\ell} \text{sgn}(\cos\theta_\ell) &= \frac{G_\mu^2 m_{b,\text{pole}}^5 |V_{ts}^* V_{tb}|^2}{48\pi^3} \left(\frac{\alpha(m_b)}{4\pi}\right)^2 (1 - \hat{s})^2 \\ &\times \left\{ -6 \text{Re} \left(\tilde{C}_{7,\text{FB}}^{\text{eff}}(\hat{s}) \tilde{C}_{10,\text{FB}}^{\text{eff}}(\hat{s})^* \right) - 3\hat{s} \text{Re} \left(\tilde{C}_{9,\text{FB}}^{\text{eff}}(\hat{s}) \tilde{C}_{10,\text{FB}}^{\text{eff}}(\hat{s})^* \right) + A_{\text{FB}}^{\text{Brems}}(\hat{s}) \right\}, \end{aligned} \quad (11)$$

and θ_ℓ denotes the angle between the momentum of the positively charged lepton and the B -meson in the rest frame of the lepton pair. The effective Wilson coefficients in Eq. (10) take into account all $O(\alpha_s)$ real and virtual corrections,^{5,9,10} as indicated by the subscript FB, whereas the last term encodes the $O(\alpha_s)$ bremsstrahlung corrections,¹² which have not been included in our analysis of q_0^2 , since they turn out to be below 1%. The NNLO QCD prediction of the normalized FB asymmetry as a function of q^2 is presented in the right plot of Figure 1, which also illustrates the numerical size of the long-distance contributions due to intermediate $c\bar{c}$ resonances.

Since q_0^2 is known to be especially sensitive to physics beyond the SM, a comprehensive study of the residual theoretical error attached to it is important. As pointed out earlier,¹⁰ the renormalization scale dependence of q_0^2 computed from Eq. (11) is rather small. However, an alternative way to estimate the residual theoretical error follows from the observation that q_0^2 should be independent of the normalization for the FB asymmetry. As both numerator and denominator of Eq. (10) are truncated series in α_s , one can expand Eq. (10) in α_s , making q_0^2 sensitive to different combinations of higher order terms, that depend on the adopted normalization. A thorough analysis shows,¹³ that the dependence of q_0^2 on the specific method used to compute it is generally larger than the scale uncertainty associated with it. Implementing all known perturbative and non-perturbative corrections, and assigning to the central value a rather

^cWe are grateful to G. Isidori for providing us with the two plots shown in Figure 1.

conservative theoretical error of 6%, that covers the whole range of possible values obtained by changing the normalization in Eq. (10), we finally get

$$q_0^2 = (3.76 \pm 0.22_{\text{theory}} \pm 0.24_{m_b}) \text{ GeV}^2, \quad (12)$$

with a total error close to 9%. Also in this case the error due to the b -quark mass can be drastically reduced by performing the calculation in a different mass scheme.

5 Outlook

Rare semileptonic B decays are going to play an important role in the search for NP at the B -factories and upcoming flavor physics experiments at the Tevatron and the LHC. From the theoretical point of view the inclusive mode in both the low- q^2 and high- q^2 window is particular interesting, since it can be accurately computed in the SM. Given this situation, a measurement of the differential decay rate integrated separately over the two reference regions would be desirable. Furthermore, as the FB asymmetry can change drastically in scenarios of NP, even a rather crude measurement of its shape would either rule out large parts of the parameter space of the underlying model or show clear evidence for physics beyond the SM.

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